# Gauge fixing of open superstring field theory in the Berkovits non-polynomial formulation

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## Abstract

We consider gauge fixing of open superstring field theory formulated by Berkovits, concentrating on the Neveu-Schwarz sector. In the free theory, we perform gauge fixing completely, which requires infinitely many ghosts and antighosts carrying various world-sheet ghost numbers and picture numbers. In the interacting theory, we have determined the form of interactions cubic in fields and antifields in the Batalin-Vilkovisky formalism.

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[*] This talk is based on the work in collaboration with Nathan Berkovits, Michael Kroyter, Yuji Okawa, Martin Schnabl and Barton Zwiebach.

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## §1.  Introduction

In string perturbation theory, closed strings appear in loop diagrams of open strings. However, in the framework of string field theory (SFT), it is non-trivial whether and how closed strings can be described in terms of open string fields. Therefore, we would like to know if open SFT can be consistently quantized without additional degrees of freedom, such as closed string fields.

In bosonic SFT,[1] quantization was discussed[2] by using the Batalin-Vilkovisky (BV) formalism,[3] but there was a difficulty caused by divergences from tadpole diagrams. Moreover, the theory also has difficulty with tachyons: quantization can be considered only in a formal way. By contrast, in superstring field theory (SSFT), we expect that these difficulties will be absent. Recent analytic methods in classical bosonic SFT[1] which developed mainly after the construction of the Schnabl solution[4] have been applied to SSFT, and the work 5) by Kiermaier and Zwiebach opened up a vista of extending the methods to at least the one-loop level. Now we are at the stage for considering quantization of open SSFT seriously.

Various SSFTs have been proposed since the work by Witten.[6] One approach to formulating an SSFT is using the picture-changing operators[7] (PCOs). At the conference, talks on gauge fixing of the SSFTs of this type, 8) and 9), were presented by M. Murata and M. Kroyter, respectively.[10] Another approach, which is achieved by Berkovits,[11] is to construct a theory in the large Hilbert space without using any PCOs. As a step toward quantizing this SSFT, we will first gauge-fix the theory by the BV formalism: we solve the *master equation*, which is a sort of Ward-Takahashi identity, and then impose gauge-fixing conditions on the solution. In the following sections, concentrating on the Neveu-Schwarz (NS) sector, we determine gauge-fixing conditions and solve the equation up to cubic order in fields and antifields.

## §2.  Gauge Fixing of the Free Theory

The free NS-sector action and its gauge symmetry of the SSFT 11) are given by[*]

$$S_0^{\text{free}} = -\frac{1}{2} \int \Phi_{(0,0)}(Q\eta_0\Phi_{(0,0)}) = -\frac{1}{2}\langle\Phi_{(0,0)}|Q\eta_0|\Phi_{(0,0)}\rangle \,, \tag{2·1a}$$

$$\delta_0\Phi_{(0,0)} = Q\epsilon_{(-1,0)} + \eta_0\epsilon_{(-1,1)} \,. \tag{2·1b}$$

Here $\Phi_{(0,0)}$ is a Grassmann-even NS-sector string field, $Q$ is the BRST operator in the first-quantized theory, and $\eta_0$ is the zero mode of $\eta$, which appears in the bosonization of the superconformal ghosts: $\beta \cong \mathrm{e}^{-\phi}\partial\xi$, $\gamma \cong \eta\mathrm{e}^{\phi}$.[7] The integration and the multiplication of

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[*] We have appended the subscript "0" on the action and the gauge variation for later convenience.

string fields are given by Witten's sewing and gluing prescription.[1] The bracket $\langle \ \ | \ \ \rangle$ is the Belavin-Polyakov-Zamolodchikov (BPZ) inner product:[12] $|\Phi_{(0,0)}\rangle$ is the state corresponding to the string field $\Phi_{(0,0)}$ and $\langle \Phi_{(0,0)}|$ is its BPZ conjugate. In Eq. (2·1) and in the sequel, the subscript $(g, p)$ on a string field indicates its world-sheet ghost number $g$ and picture $p$.[*] Since the theory is formulated in the large Hilbert space, an integral of string fields vanishes unless the integrand carries $(g, p) = (2, -1)$.

In the free theory, we can easily perform gauge fixing, using the Faddeev-Popov (FP) formalism. First, we eliminate the gauge symmetries associated with $Q$ and $\eta_0$ by the conditions $b_0 |\Phi_{(0,0)}\rangle = 0$ and $\xi_0 |\Phi_{(0,0)}\rangle = 0$, respectively. The resultant FP action is

$$S_1^{\text{free}} = \Big( \langle B_{(3,-1)}| \, b_0 + \langle B_{(3,-2)}| \, \xi_0 \Big) \Big( Q \, |\Phi_{(-1,0)}\rangle + \eta_0 \, |\Phi_{(-1,1)}\rangle \Big), \tag{2·2}$$

where $\Phi_{(-1,0)}$ and $\Phi_{(-1,1)}$ are ghosts, whereas $B_{(3,-1)}$ and $B_{(3,-2)}$ are antighosts. If we use the redefined antighost $\Phi_{(2,-1)}$,

$$\langle \Phi_{(2,-1)}| := \langle B_{(3,-1)}| \, b_0 + \langle B_{(3,-2)}| \, \xi_0 \,, \tag{2·3}$$

we obtain the gauge-fixed action of the form

$$S_0^{\text{free}} + S_1^{\text{free}} = -\frac{1}{2} \, \langle \Phi_{(0,0)}|Q\eta_0|\Phi_{(0,0)}\rangle + \langle \Phi_{(2,-1)}| \, \Big( Q \, |\Phi_{(-1,0)}\rangle + \eta_0 \, |\Phi_{(-1,1)}\rangle \Big) \tag{2·4}$$

with

$$\begin{bmatrix} b_0 \\ \xi_0 \end{bmatrix} \Phi_{(0,0)} = 0 \,, \quad b_0\xi_0\Phi_{(2,-1)} = 0 \,. \tag{2·5}$$

Decomposing the string fields with respect to the zero modes $c_0$ and $\xi_0$ helps one understand that the condition (2·5) really eliminates the gauge degree of freedom (2·1b). However, gauge fixing of the theory has not been completed yet. Since $Q$ and $\eta_0$ satisfy the relation

$$Q^2 = \eta_0^2 = \{Q, \eta_0\} = 0 \,, \tag{2·6}$$

the action $S_0^{\text{free}} + S_1^{\text{free}}$ is invariant under the gauge transformation of the ghosts

$$\delta_1\Phi_{(-1,0)} = Q\epsilon_{(-2,0)} + \eta_0\epsilon_{(-2,1)} \,, \tag{2·7a}$$

$$\delta_1\Phi_{(-1,1)} = Q\epsilon_{(-2,1)} + \eta_0\epsilon_{(-2,2)} \,. \tag{2·7b}$$

To remove this symmetry, we use the gauge-fixing condition of the form

$$\begin{bmatrix} b_0 & 0 \\ \xi_0 & b_0 \\ 0 & \xi_0 \end{bmatrix} \begin{bmatrix} \Phi_{(-1,0)} \\ \Phi_{(-1,1)} \end{bmatrix} = 0 \,, \tag{2·8}$$

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[*] The quantum number $(g, p)$ of $\xi$ and that of $\eta$ are $(-1, 1)$ and $(1, -1)$, respectively.

and introduce the ghosts for ghosts $\Phi_{(-2,0)}$, $\Phi_{(-2,1)}$ and $\Phi_{(-2,2)}$. The resultant FP action is

$$
\begin{aligned}
S_2^{\text{free}} &= \langle B_{(4,-1)}| \, b_0 \Big( Q \, |\Phi_{(-2,0)}\rangle + \eta_0 \, |\Phi_{(-2,1)}\rangle \Big) \\
&\quad + \langle B_{(4,-2)}| \Big[ \xi_0 \Big( Q \, |\Phi_{(-2,0)}\rangle + \eta_0 \, |\Phi_{(-2,1)}\rangle \Big) + b_0 \Big( Q \, |\Phi_{(-2,1)}\rangle + \eta_0 \, |\Phi_{(-2,2)}\rangle \Big) \Big] \\
&\quad + \langle B_{(4,-3)}| \, \xi_0 \Big( Q \, |\Phi_{(-2,1)}\rangle + \eta_0 \, |\Phi_{(-2,2)}\rangle \Big) \\
&= \langle \Phi_{(3,-1)}| \Big( Q \, |\Phi_{(-2,0)}\rangle + \eta_0 \, |\Phi_{(-2,1)}\rangle \Big) + \langle \Phi_{(3,-2)}| \Big( Q \, |\Phi_{(-2,1)}\rangle + \eta_0 \, |\Phi_{(-2,2)}\rangle \Big) , \quad (2\cdot9)
\end{aligned}
$$

where

$$
\langle \Phi_{(3,-1)}| := \langle B_{(4,-1)}| \, b_0 + \langle B_{(4,-2)}| \, \xi_0 , \quad \langle \Phi_{(3,-2)}| := \langle B_{(4,-2)}| \, b_0 + \langle B_{(4,-3)}| \, \xi_0 , \quad (2\cdot10)
$$

$$
\begin{bmatrix} b_0 & \xi_0 \end{bmatrix} \begin{bmatrix} \Phi_{(3,-1)} \\ \Phi_{(3,-2)} \end{bmatrix} = 0 , \quad b_0\xi_0 \, |\Phi_{(3,-1)}\rangle = b_0\xi_0 \, |\Phi_{(3,-2)}\rangle = 0 . \quad (2\cdot11)
$$

The gauge-fixed action so far is

$$
\begin{aligned}
\sum_{k=0}^{2} S_k^{\text{free}} &= -\frac{1}{2} \langle \Phi_{(0,0)}|Q\eta_0|\Phi_{(0,0)}\rangle + \langle \Phi_{(2,-1)}| \Big( Q \, |\Phi_{(-1,0)}\rangle + \eta_0 \, |\Phi_{(-1,1)}\rangle \Big) \\
&\quad + \langle \Phi_{(3,-1)}| \Big( Q \, |\Phi_{(-2,0)}\rangle + \eta_0 \, |\Phi_{(-2,1)}\rangle \Big) + \langle \Phi_{(3,-2)}| \Big( Q \, |\Phi_{(-2,1)}\rangle + \eta_0 \, |\Phi_{(-2,2)}\rangle \Big) \quad (2\cdot12)
\end{aligned}
$$

with the constraints.

In this way, we can continue gauge fixing step by step and obtain the completely-gauge-fixed action

$$
S^{\text{free}} = -\frac{1}{2} \int \Phi_{(0,0)} (Q\eta_0\Phi_{(0,0)}) + \sum_{n=1}^{\infty} \sum_{m=0}^{n-1} \int \Phi_{(n+1,-m-1)} \big( Q\Phi_{(-n,m)} + \eta_0\Phi_{(-n,m+1)} \big) \quad (2\cdot13)
$$

with

$$
\underbrace{\begin{bmatrix} b_0 & & & 0 \\ \xi_0 & \ddots & & \\ & \ddots & & b_0 \\ 0 & & & \xi_0 \end{bmatrix}}_{n+1} \begin{bmatrix} \Phi_{(-n,0)} \\ \vdots \\ \vdots \\ \Phi_{(-n,n)} \end{bmatrix} = 0 \quad (n \geq 0) , \quad (2\cdot14\text{a})
$$

$$
\underbrace{\begin{bmatrix} b_0 & \xi_0 & & 0 \\ & \ddots & \ddots & \\ 0 & & b_0 & \xi_0 \end{bmatrix}}_{n-1} \begin{bmatrix} \Phi_{(n,-1)} \\ \vdots \\ \Phi_{(n,-(n-1))} \end{bmatrix} = 0 \quad (n \geq 3) , \quad (2\cdot14\text{b})
$$

$$
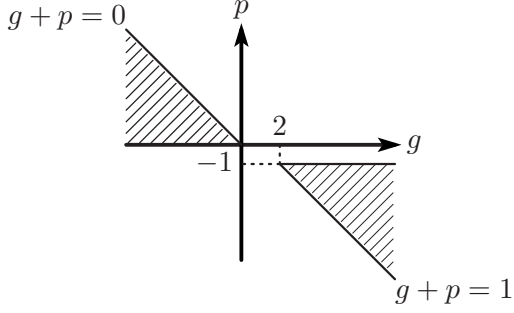b_0\xi_0\Phi_{(n,-m)} = 0 \quad (1 \leq m \leq n-1) . \quad (2\cdot14\text{c})
$$

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Fig. 1.  Distribution of allowed $(g, p)$.

Here $\Phi_{(g,p)}$ with $g \leq -1$, $0 \leq p \leq -g$ are ghosts and those with $-(g-1) \leq p \leq -1$ are antighosts. Note that a string field $\Phi_{(g,p)}$ is admissible only when the lattice point $(g, p)$ belongs to the region shown in Fig. 1: the original field $\Phi_{(0,0)}$ and the ghosts live in the left region, whereas the antighosts live in the right region in the figure.

## §3. Interacting Theory

The full NS-sector action in ref. 11) and its gauge symmetry are of the Wess-Zumino-Witten type:

$$S_0 = \frac{1}{2} \int \left( G^{-1}(QG)G^{-1}(\eta_0 G) - \int_0^1 dt \left( \hat{G}^{-1}\partial_t \hat{G} \right) \left\{ \hat{G}^{-1}(Q\hat{G}), \hat{G}^{-1}(\eta_0 \hat{G}) \right\} \right), \quad (3\cdot 1a)$$

$$\delta_0 G = \left( Q\epsilon_{(-1,0)} \right) G + G \left( \eta_0 \epsilon_{(-1,1)} \right), \quad (3\cdot 1b)$$

$$G = \exp(\Phi_{(0,0)}), \quad \hat{G} = \exp(t\Phi_{(0,0)}).$$

Note that the integral with respect to the variable $t$ is an ordinary integral. In the interacting theory, it is difficult to complete gauge fixing through the FP procedure. This is because the FP action has gauge invariance which requires the equation of motion. To see this, it is convenient to set

$$\epsilon_{(-1,0)} = e^{\phi}\tilde{\epsilon}_{(-1,0)} e^{-\phi}, \quad \phi := \Phi_{(0,0)}, \quad (3\cdot 2)$$

and redefine $\tilde{\epsilon}_{(-1,0)}$ as $\epsilon_{(-1,0)}$. This leads to

$$\delta_0 e^{\phi} = e^{\phi} \left( \tilde{Q}\epsilon_{(-1,0)} + \eta_0 \epsilon_{(-1,1)} \right), \quad (3\cdot 3)$$

where

$$\tilde{Q}X := e^{-\phi}Q\left( e^{\phi}X e^{-\phi} \right) e^{\phi} = QX + \left[ e^{-\phi}(Qe^{-\phi}), X \right]. \quad (3\cdot 4)$$

Here and in the sequel, the bracket [ , ] means the graded commutator. In the free case, the action (2·4) is exactly invariant under the transformation (2·7) because Eq. (2·6) holds. In the interacting case, however, the relation corresponding to Eq. (2·6) is

$$\tilde{Q}^2 = \eta_0^2 = 0, \quad \{\tilde{Q}, \eta_0\}X = \left[ \frac{\delta S_0}{\delta e^{\phi}} e^{\phi}, X \right], \quad (3\cdot 5)$$

so that $\{\tilde{Q}, \eta_0\}$ vanishes only when we use the equation of motion. This makes it difficult to carry out the FP procedure (or the BRST procedure). However, there is a powerful formalism

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to deal with such a complicated gauge system systematically. It is the BV formalism,[3] and we use it to consider gauge fixing of the interacting theory.

## §4.   The BV Formalism and Antifield Number Expansion

The BV formalism is an extension of the BRST formalism. In general, gauge fixing in the BV formalism is performed as follows.[*] First, one prepares ghosts, ghosts for ghosts, and so on as well as the original fields $\phi^i$ in a given action $S_0[\phi]$. Second, for each field $\Phi^A \in \{\phi^i,$ ghosts, ghosts for ghosts, ...$\}$, one introduces an additional field called an *antifield* $\Phi_A^*$. Third, starting with the action $S_0[\phi]$, one constructs the solution $S[\Phi, \Phi^*]$ to the *master equation*

$$\sum_A \frac{\partial_r S}{\partial \Phi^A} \frac{\partial_l S}{\partial \Phi_A^*} = 0 \tag{4.1}$$

under the boundary condition

$$S[\Phi, \Phi^*]|_{\Phi^*=0} = S_0[\phi] \,. \tag{4.2}$$

Here $\partial_l$ and $\partial_r$ are the left and the right derivative, respectively. Fourth, one eliminates the gauge symmetry of $S$, which originates from the symmetries of $S_0$ and ghost actions. This is achieved by imposing on the antifields conditions of the form

$$\Phi_A^* = \frac{\partial \Psi[\Phi]}{\partial \Phi^A} \,, \tag{4.3}$$

where $\Psi[\Phi]$ is a functional of $\Phi^A$'s. Finally, one obtains the completely-gauge-fixed action $S[\Phi, \partial \Psi / \partial \Phi]$. In the BV formalism, unlike in the BRST formalism, one can remove all the gauge degrees of freedom at once including those associated with ghosts, by the condition (4.3). Moreover, the BRST invariance of the gauge-fixed action is essentially equivalent to the master equation.

Let us apply this formalism to the SSFT. Note that the action $S^{\text{free}}$ (with no constraints) obtained in the free theory satisfies the master equation of the form

$$\sum_{n=0}^{\infty} \sum_{m=0}^{n} \int \frac{\delta_r S^{\text{free}}}{\delta \Phi_{(-n,m)}} \frac{\delta_l S^{\text{free}}}{\delta \Phi_{(n+2,-m-1)}} = 0 \,. \tag{4.4}$$

Actually, we can choose gauge-fixing conditions imposed on antifields such that the antifield

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[*] We use the notation where the appearance of discrete index also indicates the presence of a spacetime variable.

of $\Phi_{(0,0)}$ and those of the ghosts are identified with the antighosts[*]

$$\Phi^*_{(g,p)} \cong \Phi_{(2-g,-1-p)}, \quad (0 \le p \le -g). \tag{4.5}$$

Therefore, in what follows, we do not have to distinguish them.

In the interacting theory, it is not so easy to construct the solution to the master equation as in the free theory. However, it is known that one can solve the equation step by step, using the expansion in *antifield number*, which is defined in the present case as

$$\mathrm{afn}\left(\Phi_{(g,p)}\right) := \begin{cases} 0 & (0 \le p \le -g) \\ g-1 & (-(g-1) \le p \le -1) \end{cases}. \tag{4.6}$$

For example, the antifield numbers of $S_0^{\mathrm{free}}$, $S_1^{\mathrm{free}}$ and $S_2^{\mathrm{free}}$ in section 2 are zero, one and two, respectively. We expand the solution $S$ in antifield number,

$$S = \sum_{k=0}^{\infty} S_k, \quad \mathrm{afn}(S_k) = k. \tag{4.7}$$

Our strategy for solving the master equation in the interacting theory is as follows.

- First, we solve the master equation to some antifield number.
- Second, from the result of the first step, we infer the complete solution and confirm its validity.

In the bosonic SFT [1], the second step is easy: the complete solution can be obtained only by removing the world-sheet ghost number constraint imposed on the string field. In the present SSFT, however, this is not the case. As we can see from the result in the free theory, the original action $S_0^{\mathrm{free}}$ contains the term involving the *product* of $Q$ and $\eta_0$, whereas the FP terms do not: they consist of the *sum* of $Q$-terms and $\eta_0$-terms. Therefore there seems to be no chance that the solution will take the same form as the original action.

Starting with Eqs. (3.1a) and (3.3), we have

$$S_1 = \int \Phi_{(2,-1)} \, \mathrm{e}^{\phi} \left( \tilde{Q} \Phi_{(-1,0)} + \eta_0 \Phi_{(-1,1)} \right), \tag{4.8a}$$

$$\begin{aligned} S_2 = \int \Big[ & \Phi_{(2,-1)} \, \mathrm{e}^{\phi} \Phi_{(2,-1)} \, \mathrm{e}^{\phi} \Phi_{(-2,1)} \\ & + \Phi_{(3,-1)} \left( \tilde{Q} \Phi_{(-2,0)} + \eta_0 \Phi_{(-2,1)} + \Phi_{(-1,0)} \tilde{Q} \Phi_{(-1,0)} + \left[ \Phi_{(-1,0)}, \eta_0 \Phi_{(-1,1)} \right] \right) \\ & + \Phi_{(3,-2)} \left( \tilde{Q} \Phi_{(-2,1)} + \eta_0 \Phi_{(-2,2)} + \Phi_{(-1,1)} \eta_0 \Phi_{(-1,1)} \right) \Big]. \end{aligned} \tag{4.8b}$$

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[*] Strictly speaking, under the identification (4.5), Eq. (4.4) corresponds to the master equation for the *minimal set* of fields, which consists of the original fields and ghosts. Not until the *non-minimal* set is introduced and Lagrange multiplier fields are integrated out, is the gauge-fixing condition (2.14) imposed. For detail, see ref. 13), for example.

We have calculated $S_3$ and $S_4$ as well, but have not been able to infer a complete solution so far. For cubic interactions, however, a complete form has been obtained, as we will show in the next section.

## §5.  Cubic Interactions

It is known that the solution to the classical master equation is unique only up to *canonical transformations*, which correspond to a subset of the degrees of freedom of field redefinition. We expect that an appropriate field redefinition will help us infer a solution. In order to find such a field redefinition, we use the $\mathbb{Z}_2$-transformation property of the original action (3·1a): the action $S_0$ becomes $-S_0$ under the $\mathbb{Z}_2$-transformation

$$\left(\Phi_{(0,0)}, Q, \eta_0\right) \rightarrow \left(-\Phi_{(0,0)}, \eta_0, Q\right). \tag{5·1}$$

Let us extend this transformation to all the ghosts and antighosts obtained in the free case. We readily find that $S^{\text{free}}$ is transformed into $-S^{\text{free}}$ under

$$\left(\Phi_{(g,p)}, \Phi_{(2-g,-1-p)}, Q, \eta_0\right) \rightarrow \left(-\Phi_{(g,-g-p)}, +\Phi_{(2-g,g+p-1)}, \eta_0, Q\right) \quad (0 \leq p \leq -g). \tag{5·2}$$

Using the antifield number expansion, we have constructed the solution up to cubic order in fields and antifields which respects this extended $\mathbb{Z}_2$-transformation property. If we write the coupling constant $g$ explicitly, the result takes the form

$$S = S^{\text{free}} + g\, S^{\text{cubic}} + O(g^2), \tag{5·3}$$

$$S^{\text{cubic}} = \int \left[ -\frac{1}{6}\phi\left[Q\phi, \eta\phi\right] + \Phi^*\left(\frac{1}{2}\left[Q\phi, \Phi\right] - \frac{1}{2}\left[\eta\phi, \Phi\right]\right) + \Phi^*\Phi^*\Phi \right.$$
$$+ \frac{1}{2}\sum_{n=1}^{\infty}\sum_{m=n+2}^{\infty}\sum_{k=1}^{m-n-1}\left(\left[\Phi_{(m,-k)}, \Phi_{(1+n-m,-1+k)}\right]\eta\Phi_{(-n,1)}\right.$$
$$\left.\left. - \left[\Phi_{(m,-m+k)}, \Phi_{(1+n-m,-n+m-k)}\right]Q\Phi_{(-n,n-1)}\right)\right]. \tag{5·4}$$

Here $\Phi$ and $\Phi^*$ are the sum of all the ghosts and the sum of all the antighosts, respectively,

$$\Phi := \sum_{n=1}^{\infty}\sum_{m=0}^{n}\Phi_{(-n,m)}, \quad \Phi^* := \sum_{n=2}^{\infty}\sum_{m=1}^{n-1}\Phi_{(n,-m)}. \tag{5·5}$$

We can easily confirm that this action $S$ is transformed into $-S$ under the extended $\mathbb{Z}_2$-transformation (5·2). Moreover, $S$ is really a solution to the master equation in the following sense:

$$\sum_{n=0}^{\infty}\sum_{m=0}^{n}\int \frac{\delta_r S}{\delta\Phi_{(-n,m)}} \frac{\delta_l S}{\delta\Phi_{(n+2,-m-1)}} = O(g^2). \tag{5·6}$$

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If we impose on $S$ the gauge-fixing conditions given in section 2, we obtain the gauge-fixed action at this order.

We have now determined the cubic interactions in the NS sector completely. If we have those in the Ramond sector as well, we can extend the analysis 14) by Ellwood, Shelton and Taylor to the SSFT. Then it would be exciting to see if the gauge invariance is preserved or anomalous for one-point functions at one loop, which contain closed strings.

## Acknowledgements

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